

Natural $h \rightarrow 4g$ in Supersymmetric Models and R -Hadrons at the LHC

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Abstract

We construct a simple and natural supersymmetric model where the dominant Higgs decay is $h \rightarrow aa$ followed by $a \rightarrow gg$. In this case $m_h < m_Z$ is compatible with all experimental searches, completely eliminating the fine tuning otherwise required to satisfy Higgs search limits. The model extends the MSSM with singlet Higgs fields as well as vector-like colored particles that mediate the decay $a \rightarrow gg$. The a is a pseudo-Nambu Goldstone boson of a new global $U(1)$ symmetry, and can naturally have any mass from a few GeV to $m_h/2$. All interactions can be perturbative up to the GUT scale, and gauge coupling unification is preserved if the colored mediators come in complete GUT representations. In this case $a \rightarrow \gamma\gamma$ has a $\sim 1\%$ branching ratio, so $h \rightarrow gg\gamma\gamma$ may be observable. The colored particles that mediate the $a \rightarrow gg$ decay must be below the TeV scale, and can therefore be produced at the LHC. If these particles are stable on collider timescales, they will appear as R -hadrons, a signal visible in early LHC running. A smoking-gun signal that the stable colored particles are mediators of $h \rightarrow 4j$ is R -hadron production in association with an a . We show that this signal with $a \rightarrow \gamma\gamma$ is observable at LHC with as little as 10 fb^{-1} of integrated luminosity. Observation of R -hadrons plus missing energy would show that the superpartner of the R -hadron is R -parity odd, and therefore not an ordinary quark or gluon.

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I. INTRODUCTION

Supersymmetry (SUSY) is a compelling framework for addressing the large hierarchy between the weak scale and Planck scale. Although there is no direct experimental evidence for the existence of superpartners, there are several indirect indications that SUSY is correct. First, the minimal supersymmetric standard model (MSSM) automatically predicts precise gauge coupling unification [1]. Second, the Higgs boson is naturally light in SUSY, so precision electroweak constraints are automatically satisfied. However, the Higgs boson is generically *too* light in the MSSM: at tree level $m_h \leq m_Z$, while LEP searches give a bound $m_h > 114$ GeV. Loop corrections can increase the Higgs boson mass only at the price of fine tuning. Since the Higgs vacuum expectation value (VEV) is fixed, the physical Higgs mass can be increased only by increasing the quartic coupling. In the MSSM this is accomplished by a heavy stop mass, which gives a contribution to the quartic coupling of order

$$\Delta\lambda \sim \frac{N_c y_t^4}{16\pi^2} \ln \frac{m_{\tilde{t}}}{m_t}. \quad (1)$$

This grows logarithmically with $m_{\tilde{t}}$. For large $m_{\tilde{t}}$ there is also a contribution to the Higgs mass term that grows quadratically with $m_{\tilde{t}}$:

$$\Delta m_H^2 \sim \frac{N_c y_t^2}{16\pi^2} m_{\tilde{t}}^2. \quad (2)$$

This large contribution must be tuned away to give the observed value of the Higgs VEV, precisely the tuning problem that SUSY is supposed to solve. Satisfying the LEP Higgs mass bound necessitates a fine tuning of order 1%, and the tuning increases exponentially with the Higgs mass.

There are a number of approaches to addressing this problem. One is to extend the MSSM to get additional contributions to the quartic coupling that are not fine tuned [2–5]. Another approach is to extend the MSSM so that the Higgs decays in a non-standard way, weakening the experimental limit $m_h > 114$ GeV [6–11]. Within the MSSM, Ref. [12] argues that a specific region of parameter space with large A terms has small fine tuning, while Ref. [13] argues that anthropic considerations may explain the fine tuning. Some have advocated a large coefficient for the $SH_u H_d$ term of the NMSSM as a remedy for the fine-tuning[14]. Indeed, this can give an appreciable contribution to the Higgs boson mass but at the expense of a Landau Pole at a low scale. However, see [15–18] for approaches to make

this consistent with unification. But the most popular approach by far is simply to ignore the problem and study the fine-tuned MSSM.

In this paper we take this naturalness problem seriously and further investigate non-standard Higgs decays as a possible solution. Searches for many non-standard Higgs decays have been performed, and many are almost as sensitive as the search for Standard Model Higgs (see Ref. [19] for a review). We will focus on the cascade decay channel $h \rightarrow aa \rightarrow 4j$, which has significantly weaker limits than standard Higgs decays. The strongest published limits on this decay come from the OPAL experiment at LEP [20]. The search was designed for $h \rightarrow jj$ and the jets from a light a decay often mimic a single jet, so as a result these bounds exclude only light a masses, $m_a \lesssim 10$ GeV for $m_h < 86$ GeV. There is also a model-independent limit $m_h > 82$ GeV from OPAL for Zh production, looking for the Z recoiling against an arbitrary final state [21]. It is likely that a dedicated $h \rightarrow 4j$ search at LEP will give stronger constraints. However, no published result exists, and it is unclear whether a small excess in this channel would have been noticed.

Is it natural for $h \rightarrow aa \rightarrow 4j$ to dominate? Since $y_b^2 \sim 10^{-3}$, it is not difficult to construct models where another 2-body channel such as $h \rightarrow aa$ dominates over $h \rightarrow \bar{b}b$. However, ensuring that $a \rightarrow jj$ is the dominant a decay is more challenging. Decays to quarks require flavor-violating couplings of the a , and the decay to the heaviest quark generally dominates. Such a scenario does not significantly reduce the experimentally allowed Higgs mass since there are strong limits from LEP on $h \rightarrow 4b$ [22]. We therefore focus on the possibility that $a \rightarrow gg$ dominates. This decay can arise from the non-renormalizable operator

$$\Delta\mathcal{L} = \frac{1}{\Lambda} a \tilde{G}_{\mu\nu} G^{\mu\nu}, \quad (3)$$

where we assume that a is a pseudoscalar. This coupling can be generated by a loop diagram involving colored fields with a Yukawa coupling to a . The result is that the partial width is suppressed by both a loop factor and a heavy scale. This suppression is potentially problematic since the a field can mix with the pseudoscalar Higgs boson of the MSSM, and so the suppressed decay $a \rightarrow jj$ must compete with the mixing induced $a \rightarrow \bar{b}b$. Because of this, the simplest models will not give rise to $h \rightarrow 4g$. For example, the NMSSM with the addition of the interaction Eq. (3) will not give the desired phenomenology. In that model $m_a < m_h/2$ only near the R-symmetric or Peccei-Quinn symmetric limits. In both cases the a lives partially in the SM Higgs fields, and so the tree level decay to quarks dominates over

the loop suppressed decays to gluons.

There are already SUSY models in the literature where $h \rightarrow 4g$ dominates. The pioneering work is Ref. [8], in which a is the pseudoscalar in a gauge singlet superfield S . While this model is technically natural, allows m_a up to $m_h/2$, and represents a proof-of-principle, it has some undesirable features. For example, it has a UV divergent tadpole for S and requires non-standard soft SUSY breaking terms. Ref. [10] constructed a SUSY little Higgs theory in which a is a pseudo Nambu-Goldstone boson (PNGB). In this model $a \rightarrow gg$ dominates only if $a \rightarrow \bar{b}b$ is kinematically forbidden. This model is rather elaborate; it requires an extension of the Standard Model gauge group at the TeV scale, new flavor-dependent couplings, and large Yukawa couplings with Landau poles not far above the TeV scale.

In this paper we construct a simple and natural model in which $h \rightarrow aa \rightarrow 4g$ dominates. Our model extends the Higgs sector with gauge singlet superfields, and the decay $a \rightarrow gg$ is mediated by additional vector-like colored fields. The a is the PNGB of an approximate $U(1)$ global symmetry, and can therefore be naturally light. The decay $a \rightarrow gg$ naturally dominates over $a \rightarrow \bar{b}b$ because the latter is automatically suppressed by additional powers of explicit $U(1)$ breaking. UV divergent tadpoles and mixing terms for the singlets are forbidden by symmetries. The model works for $\text{GeV} \lesssim m_a \lesssim \frac{1}{2}m_h$, motivating experimental searches over the full kinematically allowed range. Furthermore, our model has none of the undesirable features of the previous models for $h \rightarrow 4g$. It is compatible both with grand unification and standard soft SUSY breaking terms. The model also has no dimensionful SUSY invariant couplings, and therefore preserves the solution of the μ problem of the next-to-minimal supersymmetric standard model.

This paper also points out a possible “smoking-gun” signature of models in which the decay $h \rightarrow aa \rightarrow 4g$ dominates. The signal arises from the production of the colored particles X that mediate the $a \rightarrow gg$ decay. The TeV scale is the natural scale of this theory, so the X will be copiously produced at the LHC. The X particles can decay only via flavor-dependent couplings. These must be highly suppressed because of flavor bounds, motivating (but not requiring) that they are stable on collider scales. If this is the case they will appear as R -hadrons at the LHC. This is a possible early signal at the LHC, but it is certainly not unique to our model (see [23] for a review and list of references). The novel observation made here is that if R -hadrons arise from mediators of $a \rightarrow gg$, then there is a significant cross section for a production in association with the R -hadrons, *i.e.* $X\bar{X}a$

final states. Unfortunately, $a \rightarrow gg$ is probably not observable at the LHC in such events. However, gauge coupling unification suggests that the new colored fields are embedded in GUT multiplets, in which case they will be electrically charged. This gives a branching ratio for $a \rightarrow \gamma\gamma$ of order 1%. One can then search for $X\bar{X}a \rightarrow X\bar{X}\gamma\gamma$. This may be observed with as little as 10 fb^{-1} at the LHC. This signal directly probes the $X\bar{X}a$ coupling, and together with the non-observation of a standard model Higgs boson would provide strong evidence that the R -hadrons arise from colored particles that mediate exotic Higgs decays.

II. A PNGB MODEL

The model extends the MSSM with gauge singlet superfields S , N , and \bar{N} , with an approximate global $U(1)$ symmetry under which N and \bar{N} have opposite charge. In the $U(1)$ symmetry limit, the a to which the Higgs boson decays is contained in the N and \bar{N} fields (see Eq. (7) below). In addition, there is a vector-like pair of colored superfields X and \bar{X} . The $U(1)$ invariant terms in the superpotential are

$$W = \lambda_H S H_u H_d + \frac{\kappa_S}{3} S^3 + \lambda_N S \bar{N} N + y_X N \bar{X} X. \quad (4)$$

This is $U(1)$ invariant if $\bar{X}X$ carries $U(1)$ charge. The global $U(1)$ symmetry is broken explicitly down to Z_3 by naturally small superpotential terms

$$\Delta W = \frac{\kappa_N}{3} N^3 + \frac{\kappa_{\bar{N}}}{3} \bar{N}^3. \quad (5)$$

There are actually two unbroken Z_3 symmetries. The first is a subgroup of the global $U(1)$ that acts only on N and \bar{N} , and the second is one in which all fields are re-phased by $e^{2i\pi/3}$. These symmetries are preserved by all standard soft SUSY breaking terms (scalar and gaugino masses and A terms). They forbid UV divergent tadpoles for the singlet fields, as well as UV divergent kinetic mixing among them. This is important because these effects generally make $a \rightarrow \bar{b}b$ dominate over $a \rightarrow gg$. If these Z_3 symmetries are exact, the theory has cosmologically dangerous domain walls, but very small explicit breaking is sufficient to eliminate this problem [24]. The Higgs fields all have nonzero VEVs, spontaneously breaking the approximate global $U(1)$ and giving a mass to the colored fields X and \bar{X} . For $\kappa_N, \kappa_{\bar{N}} \ll 1$, one of the pseudoscalar fields is a light PNGB which is the a particle to which the Higgs decays.

In this model the Higgs decay $h \rightarrow aa$ can easily dominate over $h \rightarrow \bar{b}b$. We first give a discussion of this point in the limit where explicit breaking of the global $U(1)$ vanishes. It is convenient to parametrize the neutral Higgs fields by

$$H_u = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ (v + h_v)s_\beta + (H_v + iA_v)c_\beta \end{pmatrix}, \quad H_d = \frac{1}{\sqrt{2}} \begin{pmatrix} (v + h_v)c_\beta - (H_v - iA_v)s_\beta \\ 0 \end{pmatrix}, \quad (6)$$

where $s_\beta = \sin \beta$ *etc.* (see *e.g.* [25]). Here h_v is not a mass eigenstate, but it is the state that unitarizes WW scattering at high energies. In the absence of fine tuning we expect one of the mass eigenstates h to be mostly h_v , and we follow common practice by calling this “the” Higgs boson. The excitations of N and \bar{N} are conveniently parametrized by

$$N = \frac{1}{\sqrt{2}}(v_N + n)e^{i(A_n + a_n)/f_N}, \quad \bar{N} = \frac{1}{\sqrt{2}}(v_{\bar{N}} + \bar{n})e^{i(A_n - a_n)/f_n}, \quad (7)$$

with $f_n = \sqrt{v_N^2 + v_{\bar{N}}^2}$. In this parametrization the a_n is derivatively coupled in the $U(1)$ symmetry limit. It also does not mix with any other field in this limit, and can therefore be identified with the massless NGB to which the Higgs decays.

The relevant coupling for $h \rightarrow aa$ decay comes from the kinetic terms for N and \bar{N} :

$$\Delta\mathcal{L} = \frac{1}{f_n^2}(v_N n + v_{\bar{N}} \bar{n})\partial^\mu a_n \partial_\mu a_n + \dots \quad (8)$$

When the n and \bar{n} scalars mix with h_v , this term can mediate the desired Higgs boson decays. The relevant decay width is given by:

$$\Gamma(h \rightarrow aa) = \frac{m_h^3(v_N U_{nh} + v_{\bar{N}} U_{\bar{n}h})^2}{32\pi f_n^4} \left(1 - \frac{2m_a^2}{m_h^2}\right)^2 \left(1 - \frac{4m_a^2}{m_h^2}\right)^{1/2}, \quad (9)$$

where U_{nh} and $U_{\bar{n}h}$ are the mixings between the Higgs boson (the mass eigenstate with the largest overlap with h_v) and the n and \bar{n} fields.

To get a concrete estimate for the branching ratio $h \rightarrow aa$, it is useful to define the combination

$$n_+ = \frac{1}{\sqrt{2}}(n + \bar{n}). \quad (10)$$

In the limit that $v_N \simeq v_{\bar{N}}$ it is this combination that appears in the relevant coupling in Eq. (8). It is this same combination that mixings with the Higgs boson via the $|F_S|^2$ term in the scalar potential:

$$V_F = \frac{1}{\sqrt{2}}\lambda_H\lambda_N v \sin(2\beta)v_N n_+ h + \dots \quad (11)$$

The mixing angle between the n_+ and the Higgs can then be estimated by dividing this result by the largest diagonal entry in the 2×2 mass matrix for n_+ and h , which we assume to be $\lambda_N v_N$. In the limit $v_N \sim v_{\bar{N}} \sim f_n$ we then obtain

$$\Gamma(h \rightarrow aa) \sim \frac{1}{16\pi} \frac{m_h^3}{v_N^2} \left(\frac{\lambda_H v \sin(2\beta)}{2\lambda_N v_N} \right)^2. \quad (12)$$

For order-1 values of the couplings, $\tan \beta$ not too large, and $f_n \sim v$ this can easily dominate over $h \rightarrow \bar{b}b$.

We now discuss the a decays. In the limit where a is an exact NGB, the only coupling linear in a is a coupling to fields that are charged under the global $U(1)$, *i.e.* X and \bar{X} . This means that the agg coupling is unsuppressed in the $U(1)$ symmetry limit. Decays such as $a \rightarrow \bar{b}b$ occur at tree level due to mixing of a and the Higgs pseudoscalar A_v , but this mixing is suppressed by the small explicit $U(1)$ breaking couplings $\kappa_N, \kappa_{\bar{N}}$. This is the basic reason that $a \rightarrow gg$ can naturally dominate even though it is loop suppressed. More precisely, assuming $\kappa_N, \kappa_{\bar{N}} \sim \epsilon \ll 1$, we have

$$\Gamma(a \rightarrow gg) = \frac{9h_{X\bar{X}a}^2 b_X^2 (N_c^2 - 1)}{64\pi} \left(\frac{g_3^2}{16\pi^2} \right)^2 \frac{m_a^3}{(y_N v_N / \sqrt{2})^2}. \quad (13)$$

On the other hand

$$\Gamma(a \rightarrow \bar{b}b) \sim \frac{N_c y_b^2}{16\pi} \epsilon^2 m_a. \quad (14)$$

Assuming that all masses and VEVs are of the same order, we have $m_a^2 \sim \epsilon v^2$, and therefore

$$\frac{\Gamma(a \rightarrow \bar{b}b)}{\Gamma(a \rightarrow gg)} \sim \frac{\epsilon y_b^2}{N_c} \left(\frac{g_3^2}{16\pi^2} \right)^{-2}. \quad (15)$$

Therefore $a \rightarrow \bar{b}b$ is sub-dominant for sufficiently small ϵ . Performing the full calculation, we find choices of parameters where ϵ can be large enough that $m_a \simeq 40$ GeV, while $a \rightarrow gg$ still dominates.

We have not undertaken an exhaustive parameter scan of this model, but it is not hard to find phenomenologically acceptable benchmark points with no fine tuning. We present one such point for concreteness. The relevant values are shown in Table I. We have used the VEVs as input parameters, with the soft mass parameters as output parameters.

This point results in a mass spectrum and branching ratios as given in Table II. LEP searches constrain $\xi_h^2 BR(h \rightarrow \bar{b}b) \lesssim 0.05$ for $m_h < 90$ GeV, and $\xi_h^2 BR(h \rightarrow \bar{b}b) \lesssim 0.1 - 0.2$ for $90 \text{ GeV} < m_h < 100 \text{ GeV}$ [22], where

$$\xi_h = \frac{g_{hZZ}}{g_{hZZ}^{(\text{SM})}}. \quad (16)$$

| | | | | |
|--------------------------------|---------------------------------------|----------------------------------|-----------------------------------|--------------------|
| $\tan \beta = 6$ | $\lambda_H = 0.13$ | $\kappa_S = -0.42$ | $\lambda_N = 0.38$ | $\kappa_N = 0.001$ |
| $A_{\kappa_N} = 0 \text{ GeV}$ | $A_{\lambda_H} = -75 \text{ GeV}$ | $A_{\kappa_S} = 200 \text{ GeV}$ | $A_{\lambda_N} = -65 \text{ GeV}$ | |
| $v_s = -1220 \text{ GeV}$ | $v_N = -425 \text{ GeV}$ | $v_{\bar{N}} = -170 \text{ GeV}$ | | |
| $y_X = 1.0$ | $\tilde{m}_X^2 = (250 \text{ GeV})^2$ | | | |

TABLE I: A set of benchmark values resulting in $h \rightarrow 4g$ decays.

| | | | | |
|----------------------------|---|---|-----------------------------|-----------------------------|
| $m_{h_1} = 78 \text{ GeV}$ | $m_{h_2} = 88 \text{ GeV}$ | $m_{h_3} = 538 \text{ GeV}$ | $m_{h_4} = 593 \text{ GeV}$ | $m_{h_5} = 709 \text{ GeV}$ |
| $m_{a_1} = 17 \text{ GeV}$ | $m_{a_2} = 455 \text{ GeV}$ | $m_{a_3} = 541 \text{ GeV}$ | $m_{a_4} = 663 \text{ GeV}$ | |
| $m_X = 300 \text{ GeV}$ | | | | |
| $\xi_{h_1}^2 = 0.18$ | $\xi_{h_1}^2 BR(h_1 \rightarrow \bar{b}b) = 0.0035$ | $\xi_{h_1}^2 BR(h_1 \rightarrow aa) = 0.18$ | $BR(a \rightarrow gg) = 1$ | |
| $\xi_{h_2}^2 = 0.81$ | $\xi_{h_2}^2 BR(h_2 \rightarrow \bar{b}b) = 0.14$ | $\xi_{h_2}^2 BR(h_2 \rightarrow aa) = 0.66$ | | |

TABLE II: Tree Level mass spectrum for the benchmark of Table I.

Therefore only a very modest contribution from the stop squarks is required to give a Higgs mass above the LEP search limits for $h \rightarrow \bar{b}b$.

III. PHENOMENOLOGY

We now turn to the phenomenology of the model. After a brief discussion of superpartner and Higgs searches, we turn to signals involving the mediator fields X and \bar{X} , where we identify a new possible distinctive signature of this class of models.

The motivation for this model is that fine-tuned radiative corrections from a large stop mass are not required. Therefore, we expect all superpartners to be near their current experimental bounds, so standard SUSY searches are expected to find superpartners early at the LHC.

Discovering the Higgs is of course more difficult in models where $h \rightarrow 4g$ dominates. There are recent analyses that claim that $h \rightarrow 4g$ may be observable using jet substructure methods with 100 fb^{-1} of LHC luminosity at 14 TeV [26, 27]. These techniques were studied for $m_a < 2m_b$, but in the model we are discussing m_a can be as large as $m_h/2$ and these search strategies are not expected to be effective in this case. For m_a near $m_h/2$, one may be able to use a boosted Higgs as studied in [26], since the four gluon jets may appear as a single fat jet. The model described in this paper motivates a detailed experimental study of

the $h \rightarrow aa \rightarrow 4g$ final state over the entire allowed kinematic range.

We now turn to the phenomenology of the X mediators. They must be below the TeV scale because they get their mass from the VEV of N , and since they are colored they will be copiously produced at the LHC. For definiteness we will assume that the fermion components of X are lighter than the scalar components (*i.e.* the soft mass-squared terms for the X scalars are positive), but our main points do not depend on this assumption. The production cross section for a color triplet fermions is shown in figure 1. Color conservation implies that X fermions must decay into an odd number of quarks, so the decay necessarily violates flavor. For example, if the mediators are part of a $\mathbf{5} \oplus \bar{\mathbf{5}}$ of $SU(5)$ (to preserve gauge coupling unification) X will have the quantum numbers of a right-handed down quark. In this case the masses of these fermions come from the superpotential terms

$$\Delta W = y_X N \bar{X} X + \epsilon_i Q_i H_d X, \quad (17)$$

where $i = 1, 2, 3$ is a flavor index. We have rotated away a possible $N \bar{X} d_i^c$ term by a field redefinition of the fields d_i^c and X . The Yukawa couplings ϵ_i mix the X with the d_i^c , allowing weak decays of the heavy mediators. Note that the dominant mass term $\bar{X} X$ preserves electroweak symmetry, and therefore does not give rise to large corrections to electroweak precision observables. The couplings ϵ_i can be made small enough to satisfy constraints from flavor and precision electroweak observables while giving prompt decays. This is a well-motivated scenario that gives rise to “fourth generation” phenomenology without unnatural tuning to satisfy experimental constraints.

However, it is also possible that Yukawa couplings such as the ϵ_i in Eq. (17) are not present or are highly suppressed. For example, they are forbidden if X is even under R -parity. In this case, the leading interaction that can give rise to X decays are higher-dimension operators such as $\Delta W \sim (\bar{X} d^c)(L H_u)$. This can easily give an X that is stable on collider scales: for example if the scale suppressing such higher-dimension operators is the GUT scale, X has a decay length of order 10^6 km.

Stable X particles will hadronize and appear in the detector as “ X -hadrons” similar to R -hadrons arising from SUSY models with gluino or squark LSP [28–32] or split supersymmetry [33]. These “ X -hadrons” may appear as highly ionizing charged tracks. The acceptance for R -hadrons at CMS is expected to be $\sim 25\%$ [34], so they can be observed in early LHC running. Early searches at CMS based on only 198 nb^{-1} of data for such exotics already

place a weak bound on this scenario. An interpolation between the results of Ref. [34] indicates a bound somewhat less than 200 GeV. This bound should improve substantially soon. Another possibility is that the X hadrons may stop in the detector and decay much later [35]. Existing searches for this signal are also sensitive to this model [36–38]. Present searches for stopped gluinos can be reinterpreted as a bound on stopping X -hadrons, with a bound of somewhat less than 300 GeV, depending on what assumptions are made about the stopping probability of the X -hadrons [37].

If R -hadrons or “fourth generation” quarks are discovered at the LHC, the next question will be what model they come from. Here we point out that a direct confirmation of the present model can come from a radiation from an X particle, since this directly probes the $a\bar{X}X$ coupling responsible for the second stage of the Higgs decay. This is given by

$$g_{a\bar{X}X} = \frac{y_X v_N}{\sqrt{2}f_n}. \quad (18)$$

The point is that discovery of either signal gives a sample of events that are essentially free of standard model backgrounds. The a decays mostly to low p_T gluon jets, and so a signature of $X\bar{X}a$ production could be $X\bar{X}jj$. However, resolving the a peak in the jj invariant mass distribution appears to be impossible because the jets have $p_T \lesssim 50$ GeV, and the energy resolution is very poor for such jets. There is a more promising signal if the X also carries electric charge, for example if X is a **5** under $SU(5)$. Then $a \rightarrow \gamma\gamma$ is also allowed, and we can search for $X\bar{X}\gamma\gamma$. If all scalars and fermions have a common mass, then

$$\text{BR}(a \rightarrow \gamma\gamma) = 3.7 \times 10^{-3}. \quad (19)$$

Backgrounds from radiation of photons or jet faking photons are negligible (as are all standard model backgrounds).

In Table III we show the production cross section for different masses of the X and a at the LHC with 14 TeV. We assume that the signal acceptance is close to that of the X -hadrons ($\simeq 25\%$) and use the branching ratio of Eq. (19). We see that we can get a handful of events in as little as 10 fb^{-1} , and a large part of the parameter space can be discovered in 100 fb^{-1} .

The heavier superpartner of the X particle (assumed here to be a scalar) also has interesting phenomenology. It decays to the lighter X particle by emitting a gluino (possibly virtual) or neutralino. If the gluino is lighter than the X partner this will give rise to a pair

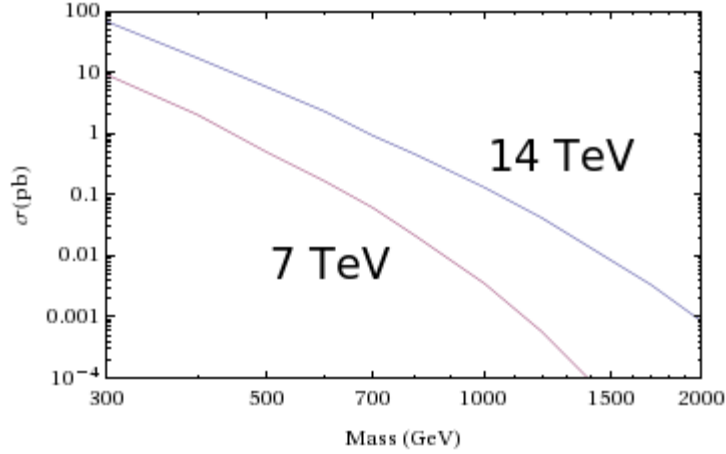


FIG. 1: Cross section for X fermion production at the LHC for different center of mass energies.

| m_X (GeV) | m_a (GeV) | $\sigma(pp \rightarrow X\bar{X}a)$ | $\text{eff} \times \sigma \times \text{BR}(a \rightarrow \gamma\gamma)$ |
|-------------|-------------|------------------------------------|---|
| 300 | 15 | 3100 | 2.9 |
| 300 | 30 | 1800 | 1.7 |
| 400 | 15 | 870 | 0.80 |
| 400 | 30 | 510 | 0.47 |
| 500 | 15 | 300 | 0.28 |
| 500 | 30 | 170 | 0.16 |
| 700 | 15 | 47 | 0.043 |
| 700 | 30 | 29 | 0.027 |

TABLE III: $\sigma(pp \rightarrow X\bar{X}a)$ in fb for different masses of the X and a , assuming the Yukawa coupling of X to a is 1, calculated using MadGraph [39]. The last column gives the expected cross section for the signal $X\bar{X}\gamma\gamma$ multiplied by a signal efficiency of 0.25 [34].

of R -hadrons together with a full SUSY cascade initiated by the gluinos. If this decay is kinematically forbidden then the decay to the LSP is expected to dominate, and we get a pair of R -hadrons plus missing energy. These striking events would directly show that the X -hadrons have a new R -parity odd superpartner. (This contrasts with R -hadrons from gluinos of squarks, where the superpartner is an ordinary particle.) In the case where the X scalar is lighter than squarks and gluinos, such processes could even be the dominant source of missing energy. Thus, dedicated searches for R -hadrons in association with missing energy

may be called for if searches for either R -hadrons or SUSY see signs of a signal.

Although their existence is not necessary for hiding the Higgs boson, unification suggests the existence of electroweak doublet partners of in the X multiplet. These might also lead to interesting collider signatures [40, 41], but the electrically charged “lepton” in the doublet is expected to rapidly decay to a soft pion and the electrically neutral “heavy neutrino”, leading to a challenging signature. With sufficient luminosity, one might observe associated production of a with these “leptons” in the final state $\gamma\gamma$ plus missing E_T , giving additional evidence for this mechanism.

IV. CONCLUSIONS

We have constructed a simple supersymmetric model in which the Higgs naturally decays dominantly via $h \rightarrow aa \rightarrow 4g$. This allows $m_h < m_Z$ and completely eliminates the need for fine tuning to satisfy the LEP Higgs bounds. Models of this kind have been considered previously in the literature, and the main difficulty is getting $a \rightarrow gg$ to dominate over $a \rightarrow \bar{f}f$ where f is the heaviest kinematically accessible fermion (generally b or τ). The main ingredient in the present model is that a is a pseudo Nambu-Goldstone boson associated with a spontaneously broken $U(1)$ symmetry in the Higgs sector. The decay $a \rightarrow \bar{f}f$ is suppressed by small explicit breaking of the $U(1)$ symmetry, while the coupling to gluons is not. Our model is significantly simpler than existing models in the literature, and it works for a masses in the full kinematically allowed range. We believe this model provides strong additional motivation for searching for the Higgs in this channel. This is especially important for larger a masses where current search strategies become ineffective.

We also pointed out a new “smoking gun” signature for models in which $h \rightarrow 4g$ is the dominant Higgs decay mode. These models necessarily have colored particles X with a large coupling to a to mediate $a \rightarrow gg$. The masses of the X particles must be below the TeV scale in order for $a \rightarrow gg$ to be large enough, so X particles can be copiously produced at LHC. They can decay only through flavor-violating couplings, and therefore may be stable on collider scales. In this case, they appear in the detector as heavy stable colored particles, “ X -hadrons.” Alternatively, if they decay, the most natural possibility is weak decays similar to a fourth generation. In either case, we expect discovery of X particles with a large number of events. This in turn gives a nearly background-free sample of X

production events, and we can look for the associated production of a with X pairs. This directly probes the coupling responsible for $a \rightarrow gg$. Associated production with $a \rightarrow gg$ is very difficult to observe, but $a \rightarrow \gamma\gamma$ is expected to have a $\sim 1\%$ branching ratio, and is readily observable. Additionally, observation of X -hadrons with missing energy is a direct sign that the X particle has a new R -parity odd superpartner.

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